Neutrino masses, mixing and leptogenesis in a two Higgs doublet model "for the third generation"

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We examine neutrino oscillations in a two Higgs doublet model (2HDM) in which the second doublet couples only to the isospin $+\frac{1}{2}$ third generation fermions. The inherently large $\tan\beta$ of this model, naturally accounts for the large top mass and, when embedded into a seesaw mechanism, also for the observed neutrino masses and mixing angles giving $-0.017 \lesssim \theta_{13} \lesssim 0.021$ at 99% CL. The large value of $\tan\beta$ also sets the mass scale of the heaviest right-handed Majorana neutrino and triggers successful leptogenesis.

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A monumental discovery of the past decade is that neutrinos have mass! This discovery is bound to have a significant impact on our understanding of the universe. It also presents us with a pressing challenge: unraveling of the mixing matrix in the lepton sector. When we recapitulate the analogous quark case we can easily grasp the importance and the difficulties of this new task. Indeed, an obvious but nonetheless crucial question is the relation of the mixing matrix in the lepton sector to that in the quark sector. Intensive experimental effort is underway and much more is being planned, round the globe, to address these issues of vital importance.

In this work we suggest that these impressive findings in the lepton sector are closely related to another remarkable finding of the 90's in the quark sector, namely that the top quark is enormously heavy compared to all the other quarks. Specifically, we will propose a simple, grounds-up approach which treats the 3rd generation neutrino in a completely analogous manner to the top quark. This reasoning leads us to a concrete framework for neutrino masses and mixings with considerable predictive power. For example, we find that the crucial mixing angle θ_{13}, which is an important target of many neutrino experiments, is quite constrained, $-0.05 < \theta_{13} < 0.035$, in our picture.

Needless to say, these observations of neutrino oscillations [1], confirming the existence of massive neutrinos, implies that the Standard Model (SM) can no longer be regarded as a minimal theory that explains all observed phenomena in Particle Physics. Bearing this in mind, our objective here is to present a simple extension of the SM, which will serve as a low energy effective theory which captures the dominant features of phenomenology and models some underlying complicated dynamics of electroweak (EW) symmetry breaking at distances much shorter than the EW scale.

Recall the allowed 3σ ranges on the atmospheric and solar neutrino square mass differences and mixing angles (including the CHOOZ $\bar{\nu}_e$ disappearance experiment) [2]:

$$\begin{aligned} 7.3 \le \Delta m^2_{sol} \cdot 10^5 \text{ eV}^2 \le 9.3 \quad &, \quad 0.28 \le \tan^2\theta_{sol} \le 0.6 \ , \\ 1.6 \le \Delta m^2_{atm} \cdot 10^3 \text{ eV}^2 \le 3.6 \quad &, \quad 0.5 \le \tan^2\theta_{atm} \le 2.1 \ , \\ \sin^2\theta_{chz} \ &\le \ 0.041 \ , \end{aligned} \tag{1}$$

where in the normal hierarchical scheme, $m_1 << m_2 << m_3$, we have $\Delta m^2_{sol} = \Delta m^2_{21}$, $\Delta m^2_{atm} = \Delta m^2_{23}$ and $\theta_{sol} = \theta_{12},\ \theta_{atm} = \theta_{23},\ \theta_{chz} = \theta_{13}$.

In order to explain the above neutrino mass differences without introducing extremely small and unnatural neutrino Dirac Yukawa couplings, the general practice is to add superheavy right handed Majorana neutrino fields, N_i, and rely on the seesaw mechanism for generating sub-eV light neutrinos:

$$\mathcal{L} = \mathcal{L}_{SM} + \mathcal{L}_Y(N, L, \phi) + M_N^{ij} N_i N_j/2 \ , \tag{2}$$

where, in the minimal setup, $\mathcal{L}_Y(N, L, \phi) = Y^\nu_{ij} L_i H N_j$, with $L_i = (\nu_{jL}, \ell_{jL})$ and H is the only Higgs doublet that couples to neutrinos. Then, the light neutrino mass matrix is given by the see-saw mechanism as:

$$m_\nu = -m_D M_N^{-1} m_D^T \ , \tag{3}$$

with $m_D = vY^\nu$ and $v =< H >$.

In this work we wish to propose a 2HDM, which we will name the "2HDM for the 3rd generation" (3g2HDM), that can naturally account for the observed neutrino masses and mixing angles, by relating the neutrino sector to the up-quark sector and giving the third generation right-handed neutrino similar status as the top-quark. This allows for a natural explanation of both the large mass of the top-quark and the apparent hierarchical structure in the neutrino sector. In particular, we extend the idea of the so called "2HDM for the top-quark" (t2HDM) [3] to the leptonic sector. In the t2HDM one assumes that ϕ_2 [the second Higgs doublet with a much larger vacuum expectation value (VEV)]

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couples *only* to the top-quark, while the other Higgs doublet ϕ_1 (with a much smaller VEV) couples to all the other fermions. The large mass hierarchy between the top quark and the other quarks becomes natural once $v_2/v_1 \equiv \tan\beta >> O(1)$, which, therefore, becomes the "working assumption" of the t2HDM.

The t2HDM has several other notable features, such as enhanced H^+bc and H^0cc couplings, new flavor changing H^0tc and H^0tu interactions and new CP phases. Thus, the t2HDM can influence CP-violation in B physics [4], flavor changing Z-decays [5] and the production of a Higgs in association with a c-jet or a b-jet in hadron colliders [6].

Following the idea of the t2HDM, we propose the following Yukawa interactions in the 3g2HDM:

$$\begin{aligned}\mathcal{L}_Y = & - Y^e \bar{L}_L \phi_1 \ell_R - Y^d \bar{Q}_L \phi_1 d_R \\ & - Y_1^u \bar{Q}_L \tilde{\phi}_1 u_R - Y_2^u \bar{Q}_L \tilde{\phi}_2 u_R \\ & - Y_1^\nu \bar{L}_L \tilde{\phi}_1 N - Y_2^\nu \bar{L}_L \tilde{\phi}_2 N \ , \end{aligned} \tag{4}$$

where N are right-handed Majorana neutrinos, Q and L are the usual quark and lepton doublets and

$$Y_1^{u,\nu} \equiv \begin{pmatrix} a^{u,\nu} & b^{u,\nu} & 0 \\ a^{u,\nu} & b^{u,\nu} & 0 \\ 0 & \delta b^{u,\nu} & 0 \end{pmatrix} \ , \ Y_2^{u,\nu} \equiv \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & c^{u,\nu} \\ 0 & 0 & c^{u,\nu} \end{pmatrix} \ , \tag{5}$$

such that, in both the quark and leptonic sectors, ϕ_2 couples only to the third generation isospin +1/2 SU(2) fermions. Note that, in (5), δ is assumed to be a real parameter of $O(1)$. Also, we have chosen $(Y_1^{u,\nu})_{13} = 0$ and $(Y_2^{u,\nu})_{31} = 0$; this is required in order to simplify our discussion below and to avoid unnatural values for some of the neutrino Yukawa couplings when confronted with the experimental limits on neutrino masses and mixings (in fact, it is sufficient to set only $(Y_2^\nu)_{13} \to 0$ in order to avoid fine tuning in the neutrino Yukawa sector). Such zero entries in Y_1 and Y_2 may result from a vanishingly small overlap between the wave functions of N and L with the corresponding indices in extra-dimensional models (see e.g., [7]), or from some flavor symmetry of the underlying short distance theory. Also, without loss of generality, we have made a further simplification by choosing similar Yukawa couplings within each generation. That is: $(Y_1^{u,\nu})_{i1}(i = 1, 2) \sim a^{u,\nu}$, $(Y_1^{u,\nu})_{i2}(i = 1, 2, 3) \sim b^{u,\nu}$ and $(Y_2^{u,\nu})_{i3}(i = 2, 3) \sim c^{u,\nu}$.

Then, $\mathcal{L}_{Y^\nu}$ yields the following Dirac neutrino mass matrix (from now on we drop the superscript ν):

$$m_D \sim \frac{v_1}{\sqrt{2}} \begin{pmatrix} a & b & 0 \\ a & b & ct_\beta \\ 0 & \delta b & ct_\beta \end{pmatrix} \ , \tag{6}$$

where $t_\beta \equiv \tan\beta = v_2/v_1$. In the basis where M_N is diagonal (one can always choose a basis for the fields N such that M_N is diagonal), $M_N = M \cdot diag(\epsilon_{M1}, \epsilon_{M2}, \epsilon_{M3})$, we then obtain from the seesaw mechanism in (3):

$$m_\nu = m_\nu^0 \begin{pmatrix} \epsilon & \epsilon & \delta\bar{\epsilon} \\ \cdot & \epsilon + \omega & \delta\bar{\epsilon} + \omega \\ \cdot & \cdot & \delta^2\bar{\epsilon} + \omega \end{pmatrix} \ , \tag{7}$$

where

$$\epsilon \equiv \frac{a^2}{\epsilon_{M1}} + \frac{b^2}{\epsilon_{M2}} \ , \ \bar{\epsilon} \equiv \epsilon - \frac{a^2}{\epsilon_{M1}} \ , \ \omega \equiv \frac{c^2 t_\beta^2}{\epsilon_{M3}} \ , \tag{8}$$

and

$$m_\nu^0 \equiv (v_1)^2/2M \ . \tag{9}$$

In the limit $\omega >> \epsilon, \ \bar{\epsilon}, \ \delta$, the physical neutrino masses and mixing angles are then given by [8]:

$$\begin{aligned} & \tan 2\theta_{23} \sim \frac{2r\omega}{\epsilon(\delta^2 - r)}, \ \tan 2\theta_{12} \sim 2\frac{g}{f}, \ \theta_{13} \sim \frac{\epsilon(\delta + r)}{2^{3/2} r\omega} \ , \\ & m_1 \sim \epsilon m_\nu^0 \left\{1 - g\sin 2\theta_{12} + f\sin^2\theta_{12}\right\} \ , \\ & m_2 \sim \epsilon m_\nu^0 \left\{1 + g\sin 2\theta_{12} + f\cos^2\theta_{12}\right\} , \\ & m_3 \sim 2\omega m_\nu^0 \ , \end{aligned} \tag{10}$$

where we have defined the parameters:

$$r \equiv \epsilon/\bar{\epsilon}, \ g \equiv \mid r - \delta \mid /\sqrt{2}/r, \ f \equiv (\delta^2 - 2\delta - r)/2/r \tag{11}$$

In the following, we will assume normal hierarchy for the light neutrinos: $m_3 >> m_2 >> m_1$, therefore leading to $m_3 \sim \sqrt{\Delta m^2_{atm}}$ and $m_2 \sim \sqrt{\Delta m^2_{sol}}$. Thus, setting $2m_\nu^0 \sim \sqrt{\Delta m^2_{sol}}$, it follows from the expression for m_3 in (10) that $\omega \sim \sqrt{\Delta m^2_{atm}/\Delta m^2_{sol}}$.

Note that, using the definitions for ω in (8), for m_ν^0 in (9) and the fact that $m_t \sim ct_\beta \times (v_1/\sqrt{2})$, our model yields the following seesaw triple-relation between the heaviest light neutrino, the heaviest right-handed Majorana neutrino and the top quark:

$$m_3 \sim 2\frac{m_t^2}{M_{N_3}} \ , \tag{12}$$

Moreover, from (9) with $2m_\nu^0 \sim \sqrt{\Delta m^2_{sol}}$ we obtain the typical mass scale (M) of the heavy right handed neutrinos:

$$M \sim \frac{v_1^2}{\sqrt{\Delta m^2_{sol}}} \sim 2\frac{m_t^2}{t_\beta^2\sqrt{\Delta m^2_{sol}}} \sim 10^{13} \text{ GeV} \ , \tag{13}$$

where, in the last equality, we have taken $m_t \sim v_2/\sqrt{2}$ and $t_\beta \sim O(10)$, which are the "working assumption" values within the 3g2HDM.

In the following numerical analysis, we will set $\omega \sim \sqrt{\Delta m^2_{atm}/\Delta m^2_{sol}} = 5.18$ and $m_\nu^0 \sim \sqrt{\Delta m^2_{sol}}/2 = 4.53 \cdot 10^{-3}$ eV, which correspond to the best fitted values [2]: $\Delta m^2_{atm} = 2.2 \cdot 10^{-3}$ [eV]^2 and $\Delta m^2_{sol} = 8.2 \cdot 10^{-5}$ [eV]^2. We note, though, that the best fitted set of input parameters (obtained by performing a minimum χ^2 fit of our model to the experimentally measured values of atmospheric and solar neutrino masses and mixing angles [2]) is:[17] $\omega \sim 5.34$, $\epsilon \sim 0.57$, $\delta \sim -1.28$ and $r \sim 1$.

In Fig. 1 we give a scatter plot of the allowed ranges in the $\theta_{13} - \theta_{23}$ and $\theta_{13} - \theta_{12}$ planes, i.e., subject to the 3σ limits listed in (1). This is done by randomly varying the three input parameters ϵ, r and δ (requiring $\omega >> \epsilon,\ \bar{\epsilon},\ \delta$) with a sample of $3 \cdot 10^6$ points. We see that the 3g2HDM predicts θ_{13} to lay within (in radians): $-0.05 \lesssim \theta_{13} \lesssim 0.035$. In fact, a minimum χ^2 analysis with respect to θ_{13} yields:

$$-0.017 \lesssim \theta_{13} \lesssim 0.021 \quad 99\%\ \mathrm{CL}\ , \tag{14}$$

with the best fitted value at $\theta_{13} \sim -0.011$. In Fig. 2

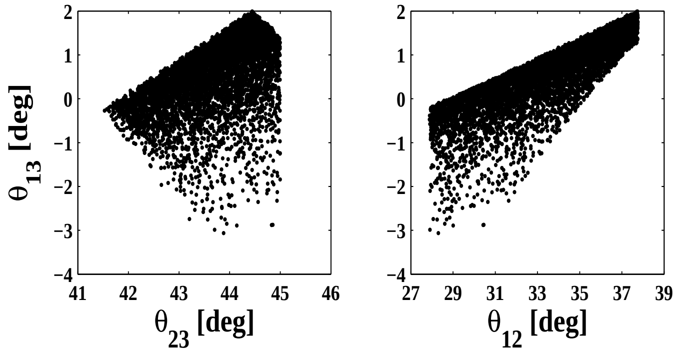

FIG. 1: *The allowed ranges in the $\theta_{13} - \theta_{23}$ (left) and $\theta_{13} - \theta_{12}$ (right) planes, for $\omega \sim 5.18$ and $m_\nu^0 \sim 4.53 \cdot 10^{-3}$ eV. See also text.*

we give a scatter plot of the allowed regions in the $\delta - \epsilon$ plane, in two limiting cases:[**?**]

case I: $r \sim 1$ ($\epsilon \sim \bar{\epsilon}$), corresponding to $\frac{a^2}{\epsilon_{M1}} << \frac{b^2}{\epsilon_{M2}}$.

case II: $r \sim 2$ ($\epsilon \to 2\bar{\epsilon}$), corresponding to $\frac{a^2}{\epsilon_{M1}} \sim \frac{b^2}{\epsilon_{M2}}$.

We see that case I (corresponding to the best fitted value for r [9]) is compatible with neutrino oscillation data for $0.4 \lesssim \epsilon \lesssim 0.7$ with $-1.7 \lesssim \delta \lesssim -1$, while the allowed range for case II is $0.5 \lesssim \epsilon \lesssim 0.8$ with $-2.4 \lesssim \delta \lesssim -1.4$. As an example, in table I we give the neutrino mixing angles and masses in cases I and II, for some specific values of the allowed parameter space.

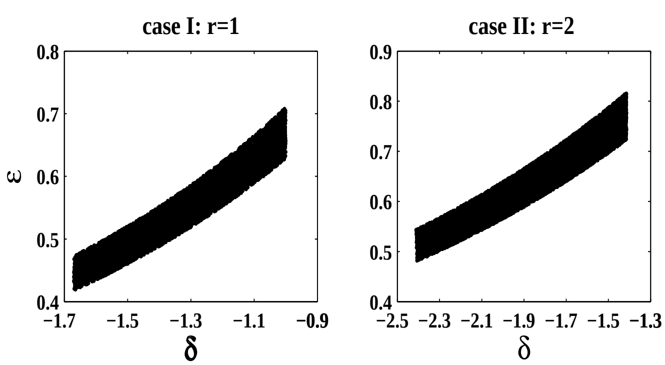

FIG. 2: *The allowed ranges in the $\epsilon - \delta$ plane, for cases I (left) and II (right). See also caption to Fig. 1.*

Let us now examine the consequences of our model on the mass spectrum of the heavy Majorana neutrinos, taking into acount the above constraints coming from neutrino oscillation. For definitness, we will consider case I

TABLE I: Mixing angles and neutrino masses in cases I and II with $\epsilon = 0.5$ and $\epsilon = 0.75$, respectively. For both cases $\delta = -1.5$, $\omega \sim 5.18$ and $m_\nu^0 \sim 4.53 \cdot 10^{-3}$ eV.

	$\delta = -1.5,\ \omega \sim 5.18,\ m_\nu^0 \sim 4.53 \cdot 10^{-3}$ eV					
case (ϵ)	θ_{23}	θ_{12}	θ_{13}	m_1 [eV]	m_2 [eV]	m_3 [eV]
I (0.5)	43^0	29.5^0	-1.14^0	0	0.0093	0.047
II (0.75)	44.7^0	35.9^0	0.69^0	$3.5 \cdot 10^{-4}$	0.0092	0.048

(i.e., $r \sim 1$, which we find as the best fitted value with respect to oscillation data), and based on quark-lepton similarity, we will take the following Ansatz for the neutrino Dirac Yukawa couplings:

$$a \sim O(10^{-3}),\ b \sim O(10^{-1}),\ c \sim O(1)\ . \tag{15}$$

Indeed, this choice follows from the assumption that the Dirac mass matrices of isospin $+1/2$ quarks and leptons exhibit a similar hierarchical structure. This Ansatz is, therefore, natural in the 3g2HDM under consideration, since this model is constructed in order to simultaneously explain the large top mass and the hierarchical structure of the neutrino masses, by relating the quark sector to the leptonic sector. That is, taking $v_1 \sim O(10)$ and $t_\beta \sim O(10)$, this Ansatz follows from the up-quark sector; $a^u v_1 \sim O(m_u)$, $b^u v_1 \sim O(m_c)$ and $m_t \sim O(c^u v_1 \tan\beta)$, if one assumes the quark-lepton similarity: $a^u \sim a^\nu \sim a$, $b^u \sim b^\nu \sim b$ and $c^u \sim c^\nu \sim c$.

It is then easy to extract the mass spectrum of the heavy Majorana neutrinos in our Ansatz, subject to the constraints given in (1). In particular, from (12) we have:

$$M_{N_3} \sim 2m_t^2/\sqrt{\Delta m^2_{atm}} \sim 10^{15}\ \mathrm{GeV}\ . \tag{16}$$

Furthermore, we find that our Ansatz (15) yields the following masses for N_2 and N_1 [for $r \sim 1$ and for $\epsilon \sim O(1)$ in accord with the constrains coming from neutrino oscillations, see Fig. 2]:

$$M_{N_2} \sim 10^{-2} M\ \ ,\ \ M_{N_1} >> 10^{-6} M\ , \tag{17}$$

which, for $M \sim 10^{13}$ GeV [see (13)], gives $M_{N_2} \sim 10^{11}$ GeV and $M_{N_1} >> 10^7$ GeV.

Let us examine how the 3g2HDM with our quark-lepton similarity Ansatz in (15) and with $r \sim 1$ and $\epsilon \sim \mathcal{O}(1)$ fits into the mechanism of leptogenesis [12]. Leptogenesis is an additional attractive feature, associated with the seesaw mechanism, since it allows the possibility of generating the observed baryon excess in the universe from a lepton asymmetry, driven by decays of the heavy Majorana neutrinos [10–12]. In particular, a CP-asymmetry, ϵ_{N_i}, in the decay $N_i \to \ell\phi_j$ can generate the lepton asymmetry (see e.g., [10, 11]):

$$n_L/s = \epsilon_{N_i} Y_{N_i}(T >> M_{N_i})\eta\ , \tag{18}$$

where $Y_{N_i} = n_{N_i}/s$, n_{N_i} being the number density of N_i and s the entropy, with $Y_{N_i}(T >> M_{N_i}) =$

$135\zeta(3)/(4\pi^4 g_*)$ and g_* being the effective number of spin-degrees of freedom in thermal equilibrium. Also, η is the "washout" parameter (efficiency factor) that measures the amount of deviation from the out-of-equilibrium condition at the time of the N_i decay.

The lepton asymmetry n_L/s can then be converted into a baryon asymmetry through nonperturbative spheleron processes. The conversion factor is [13]: $n_B/s = -(8N_G+4N_H)/(22N_G+13N_H)\times n_L/s$, where N_G and N_H are the number of generations and Higgs doublets, respectively. Thus, in our case, i.e., $N_G=3$, $N_H=2$ and $g_* \gtrsim 100$, we obtain:

$$n_B/s \sim -1.4\times 10^{-3}\epsilon_{N_i}\eta \ . \tag{19}$$

As seen from (16) and (17), our 3g2HDM can lead to a hierarchical mass spectrum for the heavy Majorana neutrinos, $M_{N_1} << M_{N_2} << M_{N_3}$, within a large portion of the allowed parameter space. In particular, imposing $M_{N_1} << M_{N_2}$, from (17) and with $M\sim 10^{13}$ GeV, we get: $M_{N_1}\sim 10^8 - 10^{10}$ GeV. With such an hierarchical mass spectrum, only the CP-asymmetry produced by the decay of N_1 survives.

In a model with two Higgs doublets, the CP-asymmetry ϵ_{N_1} is given by the sum of the individual asymmetries $\epsilon^{\Phi_1}_{N_1}$ and $\epsilon^{\Phi_2}_{N_1}$ associated with the decays $N_1\to \ell\Phi_1$ and $N_1\to\ell\Phi_2$, respectively, where Φ_1 and Φ_2 are the physical Higgs states. In particular, for $M_{N_1} << M_{N_{2,3}}$ (see e.g., [14]):

$$\epsilon^{\Phi_a}_{N_1} = -\frac{1}{16\pi(h_a^\dagger h_a)_{11}} \sum_{b=1,2}\sum_{j=2,3}\mathrm{Im}\left[(h_a^\dagger h_b)_{1j}(h_b^\dagger h_a)_{1j} + 2(h_a^\dagger h_a)_{1j}(h_b^\dagger h_b)_{1j}\right]\frac{M_{N_1}}{M_{N_j}} , \tag{20}$$

where $a=1,2$. In the 3g2HDM, the physical Higgs states Φ_1 and Φ_2 are given by the transformations: $\phi_1 = c_\beta\Phi_1 - s_\beta\Phi_2$ and $\phi_2 = s_\beta\Phi_1 + c_\beta\Phi_2$, where $s_\beta, c_\beta \equiv \sin\beta,\cos\beta$. Then, the neutrino Yukawa couplings to Φ_1 and Φ_2 are: $h_1 = c_\beta Y_1^\nu + s_\beta Y_2^\nu$ and $h_2 = -s_\beta Y_1^\nu + c_\beta Y_2^\nu$, respectively, where $Y^\nu_{1,2}$ are given in (5). Using (20), we then get (for $\tan^2\beta >> 1$):

$$\epsilon^{\Phi_1}_{N_1} = \epsilon^{\Phi_2}_{N_1} \sim -\frac{3b^2}{8\pi}\frac{M_{N_1}}{M_{N_2}}\sin 2(\theta_b-\theta_a) \ , \tag{21}$$

where we have assumed $a=|a|e^{i\theta_a}$, $b=|b|e^{i\theta_b}$ and $c=|c|e^{i\theta_c}$ in (5). Thus, using (13) and $\epsilon\sim b^2/\epsilon_{M2}$ (as implied by a fit of the 3g2HDM to neutrino oscillation data, i.e., $r\sim 1$), we obtain (setting $\epsilon_{N_1} = \epsilon^{\Phi_1}_{N_1} + \epsilon^{\Phi_2}_{N_1}$):

$$\epsilon_{N_1} \sim -\frac{3}{8\pi}\frac{t_\beta^2\sqrt{\Delta m^2_{sol}}}{m_t^2}\epsilon M_{N_1}\sin 2(\theta_b-\theta_a) \ . \tag{22}$$

The efficiency factor, η, which measures the amount of lepton asymmetry left from the CP-asymmetry ϵ_{N_1}, is calculated by solving the appropriate Boltzmann equations. The key parameter entering the Boltzmann equations is the "decay parameter", defined as [11]: $K \equiv \Gamma_{N_1}/H(T\sim M_{N_1})$, where Γ_{N_1} is the total decay width of N_1 and $H(T) = \sqrt{4\pi^3 g_*/45}T^2/M_{Plank}$ is the Hubble expansion rate. In particular, K measures the amount of asymmetry-damping caused by the inverse N_1 decay, thereby reflecting on the washout parameter η. In our model, $\Gamma_{N_1} = \Gamma^{\Phi_1}_{N_1} + \Gamma^{\Phi_2}_{N_1}$, where $\Gamma^{\Phi_{1,2}}_{N_1} = (h^\dagger_{1,2}h_{1,2})_{11}M_{N_1}/8\pi$, thus giving $\Gamma^{\Phi_2}_{N_1}/\Gamma^{\Phi_1}_{N_1} = t_\beta^2$. Therefore, for a large $\tan\beta$ (appropriate to the 3g2HDM), the decay parameter K is dominated by the decay $N_1\to\ell\Phi_2$, and is given by (for $g_*\sim 100$):

$$K \sim 4.8\times 10^{-3} a^2\cdot\frac{M_{Plank}}{M_{N_1}} \ , \tag{23}$$

which, for $a\sim 10^{-3}$ [i.e., Ansatz (15)], gives $6\lesssim K\lesssim 60$ if e.g., $M_{N_1}\sim 10^9 - 10^{10}$ GeV. In this range of values for the decay parameter (corresponding to the "mildly strong wash out" regime), the relation $\eta\sim 0.5/K^{1.2}$ constitutes a good fit to the numerical solution of the Boltzmann equations, see e.g., P. Di Bari in [11]. Thus, using the above fit for η and the CP-asymmetry in (22), the baryon asymmetry in (19) becomes:

$$\frac{n_B}{s} \sim 10^{-17}\frac{t_\beta^2\sqrt{\Delta m^2_{sol}}}{m_t^2}\epsilon M_{N_1}^{2.2}\sin 2(\theta_b-\theta_a) \ . \tag{24}$$

This has to be compared with the observed baryon to photon number ratio $n_B/n_\gamma \sim 6\times 10^{-10}$ [15], implying $n_B/s\sim 8.5\times 10^{-11}$. For example, taking $\epsilon\sim 0.5$ (see Fig. 2) and $\Delta m^2_{sol}\sim 8.2\cdot 10^{-5}$ eV^2, along with $t_\beta\sim 10$ and $m_t\sim 170$ GeV, Eq. 24 reproduces the observed baryon asymmetry for e.g., $M_{N_1}\sim 10^{10}$ GeV and $\sin 2(\theta_b-\theta_a)\sim 0.05$, or for $M_{N_1}\sim 2.7\cdot 10^9$ if CP is maximally violated in the sense that $\sin 2(\theta_b-\theta_a)\sim 1$.

Summarizing, we have presented here a simple extension to the Standard Model, in a grounds-up approach, which leads us to a framework for neutrino masses and mixings with considerable predictive power. This can clearly have significant impact on on-going as well as planned experiments. To briefly recapitulate, we have constructed a 2HDM in which the second doublet, with a much larger VEV, couples only to the third generation up-fermions (the top-quark and the 3rd generation right-handed neutrino) and the other doublet couples to all other fermions. Thus, this model envisions a close relation between quark dynamics and neutrino physics. The key parameter of this 2HDM is $\tan\beta$ which is assumed to be naturally of $O(10)$ in order to explain the large mass of the top-quark. We have shown that the large value of $\tan\beta$ in this model is directly responsible for successfully reproducing the observed neutrino oscillation data with $-0.017 < \theta_{13} < 0.021$ at 99% CL, as well as the observed baryon asymmetry of the universe through leptogenesis.

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[17] A more detailed analysis of the free parameter space of our model will be presented in a later publication [9].
[] Note that the case $\frac{a^2}{\epsilon_{M1}} >> \frac{b^2}{\epsilon_{M2}}$ is not compatible with the bounds on neutrino oscillations in (1).